

Emergent Higgs from hidden dimensions

Ryuichiro Kitano and Yuichiro Nakai

Department of Physics, Tohoku University, Sendai 980-8578, Japan

Abstract

The Higgs mechanism well describes the electroweak symmetry breaking in nature. We consider a possibility that the microscopic origin of the Higgs field is UV physics of QCD. We construct a UV complete model of a higher dimensional Yang-Mills theory as a deformation of a deconstructed (2,0) theory in six dimensions, and couple the top and bottom (s)quarks to it. We see that the Higgs fields appear as magnetic degrees of freedom. The model can naturally explain the masses of the Higgs boson and the top quark. The ρ meson-like resonances with masses such as 1 TeV are predicted.

1 Introduction

The Higgs mechanism in the Standard Model and the chiral symmetry breaking in QCD have the same structure; the nature has broken the electroweak symmetry twice in the same way. This situation motivates us to consider a possibility that the UV physics of QCD is operating as the dynamics for the electroweak symmetry breaking. An example of such a scenario has been proposed in Refs. [1, 2, 3, 4] in extra dimensional theories. If the Standard Model is defined in a compactified higher dimensional space-time, the QCD interaction gets strongly coupled at high energy and causes condensations of fermion pairs such as $\langle \bar{t}t \rangle$, breaking the electroweak symmetry. The low energy effective theory looks like the Standard Model where the Higgs boson is an emergent particle.

This scenario relies on strong gauge interactions in higher dimensional theories. However, it is this strong coupling that requires a UV cut-off to define a higher dimensional gauge theory. As in the Nambu–Jona-Lasinio model [5] for the chiral symmetry breaking, whether or not a condensation forms depends crucially on how the theory is cut-off, and thus discussion requires a UV completion of the theory.

A UV complete definition of a higher dimensional gauge theory may, in fact, be provided within four dimensional field theories by using the deconstruction technique [6, 7, 8] and supersymmetry (SUSY). It has been argued in Ref. [9] that one can obtain the (2,0) superconformal field theory in six dimensions compactified on a torus as an $N \rightarrow \infty$ limit of a four dimensional $\mathcal{N} = 2$ supersymmetric $SU(N_c)^N$ quiver gauge theory. The continuum limit provides a five dimensional $\mathcal{N} = 2$ supersymmetric gauge theory on a circle defined with two parameters: a gauge coupling g_5 and a radius R_5 . The Kaluza-Klein (KK) tower of gauge fields starts at $1/R_5$, and the theory gets strongly coupled at a scale $\Lambda_5 = 8\pi^2/g_5^2$. What is interesting is that, at this scale, it starts to appear a tower of magnetic gauge fields which can be seen from the S-duality in four dimensional supersymmetric theories. The tower represents the KK modes of a 6th dimension, and thus one can identify $\Lambda_5 = 1/R_6$. Thanks to the finiteness of the theory, we do not need to cut-off the theory at this scale even though the theory gets strongly coupled. See, for example, [10] for relations among maximally supersymmetric theories in various dimensions, and also [11, 12] for recent approaches to the (2,0) theory from five dimensional maximally supersymmetric theories.

The theory with a finite N provides the low energy effective theory of the (2,0) theory, that approximates the $N \rightarrow \infty$ limit up to a scale N/R_5 . For $1/R_6 < N/R_5$, gauge coupling constants are strong at every site, and thus one should include magnetic degrees of freedom

in the effective theory. What is important here is that, in principle, the deconstructed theory enables us to discuss physics beyond the scale Λ_5 by dealing with strongly coupled four dimensional theories.

Armed with the confidence that a higher dimensional gauge theory exists and the strong coupling simply implies the appearance of magnetic degrees of freedom, one can now consider a possibility that the Higgs field in the Standard Model is one of such emergent degrees of freedom. In this paper, based on the deconstructed (2,0) theory, we propose a two-site Standard Model as an effective theory of a higher dimensional Standard Model where the fundamental Higgs fields are absent. If one of the two QCD couplings is large, one can find that the Higgs fields appear as magnetic degrees of freedom. The four dimensional QCD interaction is kept weakly coupled since it is mainly from the other site. The electroweak symmetry breaking is described as the usual Higgs mechanism as in the Minimal Supersymmetric Standard Model (MSSM). The situation we consider is different from the models in Refs. [1, 2, 3, 4] where it is assumed that the compactification scale is of the order of $1/R_5 \sim \text{TeV}$. The scale $1/R_5$ in our framework can be much larger than TeV. The electroweak scale is rather generated by physics at $1/R_6 (\ll 1/R_5)$ where the magnetic degrees of freedom appear.

The Standard Model in a latticized extra dimension has the structure of the topcolor model [13] as noted in Ref. [14]. The two-site model we discuss is indeed similar to the super topcolor model [15], where the emergent Higgs fields are originally a pair of scalar top quarks. (See also [16] and [17] for proposals of the Higgs fields as magnetic degrees of freedom.) We will see that the Higgs boson and the top/bottom quarks obtain correct masses. Small masses compared to the naive estimates are explained by a weakly interacting feature of the magnetic theory. In the explicit model we discuss, the magnetic theory is a conformal field theory, and the Higgs/top/bottom fields interact rather weakly to the dynamical sector. This is in fact a realization of the partially composite Higgs scenario in SUSY proposed in Ref. [18].

2 Deconstructed (2,0) theory and non-supersymmetric Yang-Mills theory

It has been argued that a four dimensional $\mathcal{N} = 2$ supersymmetric $SU(N_c)^N$ gauge theory defines the (2,0) theory in six dimensions on a torus in the $N \rightarrow \infty$ limit [9]. The matter content is given in Table. 1, where $\mathcal{Q}(= Q, \bar{Q})$'s are hypermultiplets which acquire vacuum expectation values (VEVs), $Q = \bar{Q} = v\mathbf{1}$. The quiver diagram is shown in Fig. 1. The 5th dimension can be seen as the tower of the gauge fields, and the 6th one appears as the tower

	$SU(N_c)_1$	$SU(N_c)_2$	$SU(N_c)_3$	\cdots	$SU(N_c)_{N-1}$	$SU(N_c)_N$
\mathcal{Q}_1	N_c	$\overline{N_c}$	1	\cdots	1	1
\mathcal{Q}_2	1	N_c	$\overline{N_c}$	\cdots	1	1
\vdots						
\mathcal{Q}_{N-1}	1	1	1	\cdots	N_c	$\overline{N_c}$
\mathcal{Q}_N	$\overline{N_c}$	1	1	\cdots	1	N_c

Table 1: The deconstructed (2,0) theory.

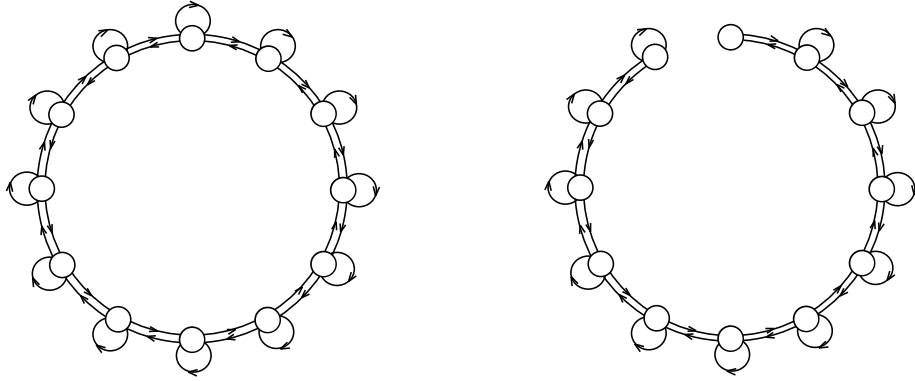


Figure 1: The quiver diagram of the deconstructed (2,0) theory (left) and an $\mathcal{N} = 1$ deformation (right).

of the magnetic states through the S-duality. The correspondence is

$$\frac{1}{R_5} = \frac{4\pi g v}{N}, \quad \frac{1}{R_6} = \frac{(4\pi)^2 v}{g} = \frac{8\pi^2}{g_5^2}, \quad (1)$$

where g is the gauge coupling constant of each site, and g_5 is the gauge coupling constant of the 5D theory. With fixed N , the quiver theory is weakly coupled when $1/R_6 > N/R_5$.

In order to take the $N \rightarrow \infty$ limit while R_5 and R_6 fixed, one needs to send $g \rightarrow \infty$. The claim of Ref. [9] is that one can take this limit since the theory is finite. One does not need to worry about the perturbativity or a possible phase transition.

One can make an IR deformation of the model to reduce to an $\mathcal{N} = 1$ SUSY Yang-Mills theory. A simple way is to add masses to \mathcal{Q}_N and an adjoint chiral superfield in the $\mathcal{N} = 2$ gauge multiplet at the 1st site, Φ_1 , such as

$$\Delta W = m Q_N \bar{Q}_N + \frac{m'}{2} \Phi_1^2. \quad (2)$$

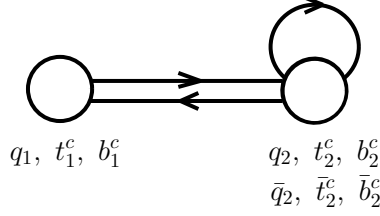


Figure 2: The quiver diagram of the two-site Standard Model.

By further adding a gaugino mass, one can obtain a non-supersymmetric Yang-Mills theory. See Fig. 1 (right) for the quiver diagram of the $\mathcal{N} = 1$ model.

The VEV v (and thus the volume) is not fixed at this level. One can fix this by gauging a baryon symmetry and add a superpotential term at the 1st site,

$$\sqrt{2}g_B Q_1 \Phi_B \bar{Q}_1 - \sqrt{2}g_B v^2 \Phi_B, \quad (3)$$

where Φ_B is the singlet chiral superfield in the $\mathcal{N} = 2$ gauge multiplet. By taking $g_B \rightarrow \infty$, the gauge multiplet becomes a Lagrange multiplier.

In general, the deconstruction technique provides a UV complete 4D theory which has the same IR physics as the extra dimensional model if the coupling constant is small. The N -site model provides the same physics as the 5D theory up to a scale N/R_5 . A lesson from the (2,0) theory is that the argument holds for strong coupling. One can keep the cut-off scale as N/R_5 whereas the strong coupling scale is set by another parameter $8\pi^2/g_5^2 = 1/R_6$ which can be lower than N/R_5 , i.e., we can discuss physics much above the energy scale $8\pi^2/g_5^2$ by including the magnetic degrees of freedom in the effective theory. Below, we couple matter fields in the Standard Model to the $\mathcal{N} = 1$ Yang-Mills model, and analyze it in a strongly coupled regime. The model is supposed to describe the IR physics of a 5D Standard Model.

3 The two-site Standard Model

We now introduce matter fields in the Standard Model in the framework. We are not aware of how we obtain such fields in a full UV finite theory. We simply introduce them and couple to the model in the previous section. Since we are only interested in the low energy physics, we consider a model with $N = 2$ which can describe physics below $1/R_5$.

As in Table. 2, we consider a model where only $SU(3)$ gauge factor has the structure of an extra dimensional theory. (See Fig. 2 for the quiver diagram.) This means that we

	$SU(3)_1$	$SU(3)_2$	$U(1)_B$	$SU(2)_L \times U(1)_Y$
Q	3	$\bar{3}$	1	1_0
\bar{Q}	$\bar{3}$	3	-1	1_0
Φ	1	$1+8$	0	1_0
q_1	3	1	1	$2_{1/6}$
t_1^c	$\bar{3}$	1	-1	$1_{-2/3}$
b_1^c	$\bar{3}$	1	-1	$1_{1/3}$
q_2	1	3	0	$2_{1/6}$
t_2^c	1	$\bar{3}$	0	$1_{-2/3}$
b_2^c	1	$\bar{3}$	0	$1_{1/3}$
\bar{q}_2	1	$\bar{3}$	0	$\bar{2}_{-1/6}$
\bar{t}_2^c	1	3	0	$1_{2/3}$
\bar{b}_2^c	1	3	0	$1_{-1/3}$

Table 2: The two-site Standard Model.

have already integrated out the KK modes of the $SU(2)_L \times U(1)_Y$ gauge fields. We are implicitly assuming that there are no significant effects at low energy from the $SU(2)_L \times U(1)_Y$ dynamics in the extra dimension. We introduce the top and bottom quark multiplets at both the 1st and the 2nd sites, representing that the matter fields are propagating in the extra dimension [19]. In order to reproduce the correct number of the chiral matter, vector-like partners are introduced at the 2nd site. For simplicity, we put other matter fields such as quarks in the 1st and 2nd generations and leptons at the 2nd site for now. Interestingly, the model is very similar to the (super) topcolor model [13, 15]. The superpotential terms relevant for the discussion are

$$W = \sqrt{2}g \left(q_1 \bar{Q} \bar{q}_2 + t_1^c Q \bar{t}_2^c + b_1^c Q \bar{b}_2^c + \bar{Q} \Phi Q - v^2 \text{Tr} \Phi + v_q \bar{q}_2 q_2 + v_t \bar{t}_2^c t_2^c + v_b \bar{b}_2^c b_2^c \right). \quad (4)$$

The mass terms at the 2nd site represent the profiles of the matter fields in the extra dimension. For

$$v_q = v_t = v_b = v, \quad (5)$$

the zero modes of the matter fields have flat distributions.

The gauge coupling constants at the 1st and the 2nd sites run differently once we deform the theory. In particular, the 1st site is asymptotically free whereas the 2nd one is IR free.

We define here the dynamical scale of $SU(3)_1$ as Λ . For

$$\Lambda \ll 4\pi v, \quad (6)$$

the classical level analysis is reliable. In this case, the low energy theory is the MSSM without the Higgs fields. The low energy QCD coupling is given by

$$\frac{1}{g_{\text{QCD}}^2} = \frac{1}{g_1^2} + \frac{1}{g_2^2}, \quad (7)$$

where g_1 and g_2 are coupling constants of $SU(3)_1$ and $SU(3)_2$, respectively.

Below we analyze the strongly coupled region $\Lambda \gg 4\pi v$. We assume that g_2 is small enough to reproduce the QCD coupling at low energy. The difference of the coupling between g_1 and g_2 may need some explanation in order for the model to be interpreted as an IR theory of a higher dimensional theory. A possibility is just a quantum effect due to the different renormalization group running. One can also assume that the gauge coupling is position dependent. For example, there can be a large localized kinetic term for the gauge boson somewhere away from the 1st site.

It is not completely clear if the model with matter fields is really a part of some UV complete higher dimensional theory, although we suspect that the string theory construction is possible. In any case, the two-site reduced model itself is a renormalizable theory and its IR behavior can be reliably analyzed as we see below. One can, of course, take the two-site model as the definition without referring to extra dimensions.

4 IR physics – the Seiberg dual picture

The theory at the 1st site is a supersymmetric QCD with $N_c = 3$ and $N_f = 5$. This theory is at the strongly coupled edge of the conformal window. For a strongly coupled region,

$$\Lambda \gg 4\pi v, \quad (8)$$

the IR physics is better described by the Seiberg dual picture [20]. The dual picture is an $SU(2)$ gauge theory with five flavors, and has a more weakly coupled fixed point. The particle content of the dual theory is given in Table. 3. The fields f , \bar{f} , f' , and \bar{f}' are dual quarks, and $H_u (\sim q_1 b_1^c)$ and $H_d (\sim q_1 t_1^c)$ are meson fields in the dual theory*. The top and bottom

*We use the convention which is commonly used in the MSSM literature, where H_u (H_d) has an hypercharge $1/2$ ($-1/2$). This is in fact a slightly confusing notation. In this model, H_u and H_d are made of the down and up-type quark superfields, respectively, and indeed, as we will see later, H_u and H_d respectively give masses to the down and up-type quarks unlike the MSSM.

	$SU(2)_1$	$SU(3)_2$	$U(1)_B$	$SU(2)_L \times U(1)_Y$
f	2	1	3/2	2_0
\bar{f}_u	$\bar{2}$	1	-3/2	$1_{1/2}$
\bar{f}_d	$\bar{2}$	1	-3/2	$1_{-1/2}$
H_u	1	1	0	$2_{1/2}$
H_d	1	1	0	$2_{-1/2}$
f'	2	3	3/2	$1_{1/6}$
\bar{f}'	$\bar{2}$	$\bar{3}$	-3/2	$1_{-1/6}$
q	1	3	0	$2_{1/6}$
t^c	1	$\bar{3}$	0	$1_{-2/3}$
b^c	1	$\bar{3}$	0	$1_{1/3}$

Table 3: The dual picture of the two-site Standard Model below the scale Λ .

quarks q , t^c , and b^c are mixtures of the mesons and those at the 2nd site. The superpotential is

$$W = h (\bar{f}_u H_d f + \bar{f}_d H_u f - \epsilon_t \bar{f}_u t^c f' - \epsilon_b \bar{f}_d b^c f' - \epsilon_q \bar{f}' q f) + \frac{(4\pi v)^2}{\Lambda} \bar{f}' f', \quad (9)$$

where

$$\epsilon_{q,t,b} = \frac{(4\pi)^2 v_{q,t,b}}{\sqrt{(h\Lambda)^2 + ((4\pi)^2 v_{q,t,b})^2}}, \quad (10)$$

and h is a dimensionless coupling constant. The scale Λ represents the dynamical scale of the $SU(3)_1$ at which massive hadronic modes appear. In general, the parameter Λ appearing in the superpotential can be an arbitrary value, and a choice would change the dynamical scale of the dual theory. We take the parameter as Λ such that the dual theory has the same dynamical scale. We put a factor of $(4\pi)^2$ from the naive dimensional analysis (NDA) [21, 22].

Since we assume $\Lambda \gg 4\pi v$, the mass term of f' and \bar{f}' is much smaller than Λ . Below the dynamical scale Λ , the theory is, therefore, approximately conformally invariant. The dimensions of the chiral superfields can be determined by the a -maximization [23]. By choosing the parameters v_q , v_t and v_b so that $\epsilon_q \sim \epsilon_t \sim 1$ and $\epsilon_b \ll 1$, which will be later required from $m_t \gg m_b$, we obtain the dimensions as follows:

$$D(H_d) = 1.03, \quad D(H_u) = 1.13, \quad D(q) = 1.13, \quad D(t^c) = 1.17, \quad (11)$$

$$D(f) = 0.99, \quad D(\bar{f}_u) = 0.99, \quad D(\bar{f}_d) = 0.88, \quad D(f') = 0.84, \quad D(\bar{f}') = 0.88. \quad (12)$$

We can see that the dimensions are close to unity, representing that the theory is weakly coupled, and thus the effective description in terms of the magnetic degrees of freedom is appropriate. By matching these dimensions to the one-loop level computations of the anomalous dimensions, we obtain the sizes of the gauge coupling and the superpotential couplings as follows:

$$\frac{\tilde{g}}{4\pi} \sim 0.41, \quad \frac{\lambda_d}{4\pi} \sim 0.11, \quad \frac{\lambda_u}{4\pi} \sim 0.26, \quad \frac{\lambda_t}{4\pi} \sim 0.29, \quad \frac{\lambda_q}{4\pi} \sim 0.26, \quad (13)$$

where \tilde{g} is the gauge coupling of $SU(2)_1$ and λ 's are coupling constants in the superpotential:

$$W = \lambda_d \bar{f}_u H_d f + \lambda_u \bar{f}_d H_u f + \lambda_t \bar{f}_u t^c f' + \lambda_b \bar{f}_d b^c f' + \lambda_q \bar{f}' q f + \frac{(4\pi v)^2}{\Lambda} \bar{f}' f', \quad (14)$$

where $\lambda_b \ll 4\pi$. Since the gauge coupling \tilde{g} is not very small, the one-loop estimation should be regarded as an order estimate. We see that λ_d is somewhat smaller than others. We keep track of these couplings λ .

The f' and \bar{f}' fields have a mass term with a coefficient $(4\pi v)^2/\Lambda$. Through the anomalous dimension of the $\bar{f}' f'$ operator, the actual mass scale Λ' is slightly larger by a relation:

$$\frac{(4\pi v)^2}{\Lambda} \left(\frac{\Lambda}{\Lambda'} \right)^{2-D(\bar{f}')-D(f')} = \Lambda'. \quad (15)$$

By integrating them out, we obtain a superpotential:

$$W = \lambda_d \bar{f}_u H_d f + \lambda_u \bar{f}_d H_u f - \frac{\lambda_q \lambda_t}{\Lambda'} f \bar{f}_u t^c q - \frac{\lambda_q \lambda_b}{\Lambda'} f \bar{f}_d b^c q. \quad (16)$$

At the mass scale Λ' , the $SU(2)_1$ factor confines due to the decoupling of the flavors. The superpotential couplings λ are now renormalized and we expect that they become larger than the fixed point values. In the following we treat λ 's as free parameters.

The low energy effective theory is described by the particles in Table. 4 with a constraint:

$$H'_u H'_d - S \bar{S} = \frac{\Lambda'^2}{(4\pi)^2}. \quad (17)$$

A factor of $(4\pi)^2$ is from the NDA. The superpotential is given by

$$W = \frac{\lambda_u \Lambda'}{4\pi} H_u H'_d + \frac{\lambda_d \Lambda'}{4\pi} H_d H'_u - \frac{\lambda_q \lambda_t}{4\pi} H'_u t^c q - \frac{\lambda_q \lambda_b}{4\pi} H'_d b^c q. \quad (18)$$

We arrive at a model similar to the MSSM.

At the supersymmetric level, there will be no electroweak symmetry breaking. The baryon fields S and \bar{S} develop VEVs by the constraint in Eq. (17). This will break the $U(1)_B$ and the

	$SU(3)_2$	$U(1)_B$	$SU(2)_L \times U(1)_Y$
H_u	1	0	$2_{1/2}$
H_d	1	0	$2_{-1/2}$
H'_u	1	0	$2_{1/2}$
H'_d	1	0	$2_{-1/2}$
S	1	3	1
\bar{S}	1	-3	1
q	3	0	$2_{1/6}$
t^c	$\bar{3}$	0	$1_{-2/3}$
b^c	$\bar{3}$	0	$1_{1/3}$

Table 4: The dual picture of the two-site Standard Model below the scale Λ' .

phase direction of S and \bar{S} gets eaten by the gauge field. If the $U(1)_B$ factor is dynamical at this scale, there should be a string associated with this symmetry breaking. Since the $U(1)_B$ factor is not originally dynamical, the string should be unstable. It implies the presence of monopoles (or dyons) at the scale Λ , which can attach to the string. For $\Lambda \gg \Lambda'$, the string is meta-stable. See [24] for a recent discussion on this string.

The massless degrees of freedom are just the MSSM without the Higgs fields. That is the same result as the case with small g , representing a smooth transition between strong ($\Lambda \gg 4\pi v$) and weak ($\Lambda \ll 4\pi v$) coupling regimes. There are various massive modes (including a string) at the scale Λ' . In particular, there is a light Higgs field H_d with a mass $\lambda_d \Lambda' / (4\pi)$. This is the emergent Higgs field responsible for the electroweak symmetry breaking as we see later.

For $\Lambda \gg 4\pi v$, the mass scale Λ' is much lower than Λ . If the model has a geometric interpretation, two physical scales Λ and Λ' can naturally be identified as

$$\frac{1}{R_5} \sim \Lambda, \quad \frac{1}{R_6} \sim \Lambda'. \quad (19)$$

If we ignore the renormalization effects, Eqs. (15) and (19) imply $v^2 R_5 R_6$ be a constant which is indeed the case in Eq. (1) for a fixed N .

For $N_c \gg 1$, one can treat the matter fields q , t^c and b^c as a probe since the dynamics is dominated by the $SU(N_c)$ gauge interaction. Also, as we take the deformation in Eq. (2) small, the Seiberg duality is suspected to smoothly connect to the S-duality in the $\mathcal{N} = 2$ SUSY gauge theory [25]. Therefore, we argue that the analysis we have done traces the

properties of the original (2,0) theory.

5 Electroweak symmetry breaking

As in the MSSM, adding soft SUSY breaking terms can create a vacuum at $H_u, H_d \neq 0$, which represents the dynamical electroweak symmetry breaking. With generic values of the Higgs VEVs, S and \bar{S} acquire VEVs and get eaten.

The SUSY breaking terms can be introduced in various ways. For example, one can explicitly break SUSY at the scale Λ' so that the effective theory below Λ' is just the non-supersymmetric Standard Model. If the scale Λ' is to be identified as the $1/R_6$, one can consider a possibility of breaking SUSY by boundary conditions of the 6th dimension as in Ref. [26]. As a similar possibility, one can introduce an F -component VEV of v and Λ which corresponds to the radion F -term, representing the SUSY breaking by the boundary condition of the 5th dimension [27, 28]. Through the relation in Eq. (15), Λ' generically acquires the F -component, and thus it is upgraded to a chiral superfield $\Lambda'(1 + m_{\text{SUSY}}\theta^2)$ where m_{SUSY} is a typical size of the SUSY breaking. With the non-zero F -component of Λ' , we obtain a structure of gauge mediation [29, 30, 31] in Eq. (9) where f' and \bar{f}' are messenger fields. For the concreteness of the discussion, we take this option of SUSY breaking in the following.

Hereafter, we assume $m_{\text{SUSY}} \sim \Lambda'$ which makes the discussion simple and, moreover, provides a successful electroweak symmetry breaking. We reserve a detailed analysis of the potential for a future work. In the following, we study the potential by treating $\lambda/4\pi$ as small parameters. All the estimates are based on the NDA where we ignore all the $O(1)$ factors including color factors. The results should be interpreted as order estimates and deviations by a factor of a few can easily happen, although all the predictions are pretty successful as we see below.

In the H_u and H_d directions, there are SUSY breaking potential such as

$$V \ni \frac{m_{\text{SUSY}}^2}{(4\pi)^2}(|\lambda_u H_u|^2 + |\lambda_d H_d|^2) + \frac{1}{(4\pi)^2}(|\lambda_u H_u|^4 + |\lambda_d H_d|^4). \quad (20)$$

Here we have omitted unknown $O(1)$ parameters. We see that the potential is suppressed by λ 's. Therefore, it is mainly the direction of the 126 GeV Higgs boson [32, 33]. This is a realization of the partially composite Higgs scenario in Ref. [18]. The H'_u and H'_d direction also has a SUSY breaking potential:

$$V \ni m_{\text{SUSY}}^2(|H'_u|^2 + |H'_d|^2) + \cdots \quad (21)$$

We assume that the quadratic term is positive and the soft terms are common for H'_u and H'_d as they are not distinguished by the $SU(2)_1$ gauge dynamics.

There are also contributions from the superpotential:

$$W \ni \frac{\Lambda'}{4\pi}(\lambda_u H_u H'_d + \lambda_d H_d H'_u) + m_{\text{SUSY}} H'_u H'_d. \quad (22)$$

The last term needs some explanation. It originates from the kinetic terms of S and \bar{S} , which can be expressed in terms of H'_u and H'_d by using the constraint in Eq. (17) as

$$K \ni 2 \left| H'_u H'_d - \frac{\Lambda'}{(4\pi)^2} \right|. \quad (23)$$

By expanding around $H'_u = H'_d = 0$, we obtain

$$K \ni \frac{\Lambda'^\dagger}{\Lambda'} H'_u H'_d + \text{h.c.} \quad (24)$$

The SUSY breaking term associated with this provides μ - and $B\mu$ -like terms

$$W \ni m_{\text{SUSY}} H'_u H'_d, \quad V \ni m_{\text{SUSY}}^2 H'_u H'_d + \text{h.c.} \quad (25)$$

These terms explicitly break anomalous $U(1)$ symmetries and thus provide masses to pseudoscalar Higgs bosons.

There are also mixing terms between H and H' fields through $B\mu$ -like terms associated with first two terms in Eq. (22):

$$V \ni m_{\text{SUSY}} \left(\frac{\lambda_u \Lambda'}{4\pi} H_u H'_d + \frac{\lambda_d \Lambda'}{4\pi} H_d H'_u + \text{h.c.} \right), \quad (26)$$

with which the size of the mixing is of order $\lambda/4\pi$. Through these mixing terms, the H' fields obtain VEVs, such as

$$\langle H'_u \rangle \sim \langle H'_d \rangle \sim \frac{\lambda_d}{4\pi} \langle H_d \rangle + \frac{\lambda_u}{4\pi} \langle H_u \rangle. \quad (27)$$

By integrating out the heavy H' fields, we obtain the effective potential of H_u and H_d as in the MSSM. In the potential $V_{\text{eff}}(H_u, H_d)$, at the leading order of $\lambda/(4\pi)$, the H_u and H_d fields are always accompanied with the coupling constants λ_u and λ_d , except for the D -terms of the MSSM gauge interactions. Moreover, there is an approximate symmetry to flip $\lambda_u H_u$ and $\lambda_d H_d$ again except for the D -terms and also from the quantum corrections from the stop loops. In this case, the potential can be written in terms of $H_+ \equiv (\lambda_d H_d + \lambda_u H_u)/\sqrt{2}$ and $H_- \equiv (\lambda_d H_d - \lambda_u H_u)/\sqrt{2}$, and that is approximately invariant under $H_\pm \rightarrow \pm H_\pm$. Therefore, the minimum can generically be at $\langle H_+ \rangle \neq 0$ and $\langle H_- \rangle \sim 0$, which we assume to be the case

motivated by the discussion of the T -parameter later. The correction from the stop loop is of order

$$\langle H_- \rangle \sim \frac{\lambda_q^2 \lambda_t^2}{(4\pi)^4} \frac{m_{\tilde{t}}^2}{m_{\text{SUSY}}^2} \langle H_+ \rangle, \quad (28)$$

where $m_{\tilde{t}}$ is the stop mass, which is of order

$$m_{\tilde{t}}^2 \sim \frac{\lambda_t^2}{(4\pi)^2} m_{\text{SUSY}}^2. \quad (29)$$

The correction from the D -term is

$$\langle H_- \rangle \sim \frac{(4\pi)^2}{\lambda_d^2} \frac{m_Z^2}{m_{\text{SUSY}}^2} \langle H_+ \rangle, \quad (30)$$

for $\lambda_u \gg \lambda_d$. Higher order corrections are of order

$$\langle H_- \rangle \sim \frac{\lambda_u^2}{(4\pi)^2} \langle H_+ \rangle, \quad (31)$$

for $\lambda_u \gg \lambda_d$. As we will see later, the VEV $\langle H_- \rangle$ should be smaller than $\langle H_+ \rangle$ by a factor of a few. We therefore assume that is the case, and concentrate on the H_+ direction, i.e., $\tan \beta = \langle H_u \rangle / \langle H_d \rangle = \lambda_d / \lambda_u$. For $\lambda_u \gg \lambda_d$, motivated by Eq. (13), the VEV direction is approximately H_d , which we assume in the following discussion.

In order to achieve the correct electroweak symmetry breaking, we need to have the correct size of the quadratic term:

$$\frac{\lambda_d^2}{(4\pi)^2} m_{\text{SUSY}}^2 \sim \frac{m_h^2}{2}, \quad (32)$$

where m_h is the observed Higgs boson mass, 126 GeV. The left-hand-side is a collection of various contributions such as from the superpotential, the Kähler potential and also from a mixing with H' . The quartic term should satisfy:

$$\frac{\lambda_d^4}{(4\pi)^2} + \frac{g_L^2 + g_Y^2}{2} \sim \frac{m_h^2}{\langle H \rangle^2} \sim 0.5, \quad (33)$$

in order to obtain the correct Higgs VEV, $\langle H \rangle = 174$ GeV. The second term is from the D -term potential of the MSSM gauge interactions, where g_L and g_Y are coupling constants of the $SU(2)_L$ and $U(1)_Y$ gauge interactions. Eq. (33) implies

$$\frac{\lambda_d}{4\pi} \sim 0.2. \quad (34)$$

The size is not far from the fixed point value in Eq. (13). The Higgs VEV should be limited by $\lambda_d \langle H_d \rangle \lesssim \Lambda'$, since otherwise the confinement scale should be redefined as $\Lambda' \sim \lambda_d \langle H_d \rangle$. Therefore,

$$\Lambda' \gtrsim 400 \text{ GeV} \cdot \left(\frac{\lambda_d/4\pi}{0.2} \right). \quad (35)$$

The relation in Eq. (32) now provides a measure for a required degree of fine-tuning in the electroweak symmetry breaking such as

$$\delta = \frac{m_h^2/2}{(\lambda_d m_{\text{SUSY}}/4\pi)^2} = 20\% \cdot \left(\frac{m_{\text{SUSY}}}{1 \text{ TeV}}\right)^{-2} \left(\frac{\lambda_d/4\pi}{0.2}\right)^{-2}. \quad (36)$$

There is essentially no fine-tuning for $m_{\text{SUSY}} \lesssim \text{TeV}$. We see that a natural electroweak symmetry breaking is achieved while $m_h = 126 \text{ GeV}$ is obtained.

From the superpotential in Eq. (22), the Higgsino masses are generated. The lightest Higgsino obtains a mass of order

$$m_{\tilde{h}} \sim \frac{\lambda_u \lambda_d}{(4\pi)^2} \frac{\Lambda'^2}{m_{\text{SUSY}}} \sim 120 \text{ GeV} \cdot \left(\frac{\lambda_d/4\pi}{0.2}\right) \left(\frac{\lambda_u/4\pi}{0.6}\right) \left(\frac{\Lambda'}{1 \text{ TeV}}\right)^2 \left(\frac{m_{\text{SUSY}}}{1 \text{ TeV}}\right)^{-1}. \quad (37)$$

In the case where the gauginos are much heavier than the Higgsinos, the charged Higgsino and the neutral Higgsino degenerate. The Higgsino mass bounds in such a case are $m_{\tilde{h}} \gtrsim 90 \text{ GeV}$, obtained from the searches for mono-photon signals at the LEP-II experiments [34, 35]. Although this is an order estimate, one can expect a quite light Higgsino if there is no significant fine-tuning in the electroweak symmetry breaking.

The top quark mass is generated by two sources. One is from the Yukawa interaction in the superpotential in Eq. (18) together with the VEV of H'_u in Eq. (27). There is also a contribution from a Kähler term:

$$K \ni \frac{\lambda_q \lambda_t \lambda_d}{(4\pi)^2} \frac{1}{\Lambda'^{\dagger}} H_d^{\dagger} t^c q, \quad (38)$$

which generates the non-holomorphic Yukawa term through SUSY breaking. Both contributions are of the order of

$$m_t \sim \frac{\lambda_q \lambda_t \lambda_d}{(4\pi)^2} \langle H_d \rangle \sim 160 \text{ GeV} \cdot \left(\frac{\lambda_d/4\pi}{0.2}\right) \left(\frac{\lambda_q/4\pi}{0.6}\right) \left(\frac{\lambda_t/4\pi}{0.6}\right). \quad (39)$$

The top quark mass, $m_t = 173 \text{ GeV}$, can be successfully explained. The bottom quark mass is also generated in the same way. The correct size can be obtained by choosing the value of ϵ_b appropriately.

It is interesting that the top quark is not required to be very heavy unlike the top condensation models [36, 37, 13]. The main Higgs directions, H_u and H_d , are originally composite of the stop and anti-stop (or sbottom and anti-sbottom). The fermion pairs such as the $t\bar{t}$ and $b\bar{b}$ condensations correspond to the F -component of H_u and H_d . From the equations of motion, $F_{H_u}^{\dagger} = -\partial W/\partial H_u$ and $F_{H_d}^{\dagger} = -\partial W/\partial H_d$, we can see that they correspond to H'_d and H'_u . The electroweak symmetry breaking in this model is mainly by

the stop condensation, and the $t\bar{t}$ condensation is a minor contribution. This explains why the top quark can be light unlike non-SUSY models.

Superparticles except for Higgsinos can be much heavier than the electroweak scale. The stop and sbottom masses are

$$m_{\tilde{t}} \sim m_{\tilde{b}} \sim \frac{\lambda_q}{4\pi} m_{\text{SUSY}} \sim 600 \text{ GeV} \cdot \left(\frac{\lambda_q/4\pi}{0.6} \right) \left(\frac{m_{\text{SUSY}}}{1 \text{ TeV}} \right). \quad (40)$$

For a light Higgsino, the stop (sbottom) can decay into $t\chi^0$ ($b\chi^0$) or $b\chi^+$ ($t\chi^-$). At the LHC experiments, chargino decays are invisible due to a small mass splitting. In this case, the search for two b -jets and missing transverse momentum put the severest constraint. The current lower bound on the sbottom mass is about 600 GeV by assuming a 100% branching fraction into $b\chi^0$ [38, 39]. Our prediction is close to this bound.

Unlike the models based on the SUSY desert, we do not expect a large logarithmic enhancement in the quantum corrections of the SUSY breaking terms. In that case, the gauginos and other squarks (including the right-handed sbottom) can be as heavy as 2 TeV while maintaining naturalness of the electroweak symmetry breaking [40, 41]. Such a spectrum is obtained in a scenario where SUSY is broken somewhere away from the 1st site.

6 Fermion masses

At this stage, the quarks in the first two generations and leptons are massless. A possible way to couple the fermions to the Higgs fields is to let them propagate into the extra dimension and write down superpotential terms such as

$$\frac{q_1 b_1^c F F_u^c}{\Lambda_F} + \frac{q_1 t_1^c F F_d^c}{\Lambda_F} \quad (41)$$

at the 1st site, where F , F_u^c and F_d^c are quark and lepton superfields. Λ_F is a mass scale where these terms are generated. These correspond to the Yukawa interaction terms at low energy:

$$\frac{\lambda_u \Lambda}{(4\pi)^2 \Lambda_F} \left(\frac{\Lambda}{\Lambda'} \right)^{1-D(H_u)} H_u F F_u^c + \frac{\lambda_d \Lambda}{(4\pi)^2 \Lambda_F} \left(\frac{\Lambda}{\Lambda'} \right)^{1-D(H_d)} H_d F F_d^c \quad (42)$$

By appropriately choosing the mixing factors, one can obtain a hierarchy of the fermion masses naturally. Note here that the scale Λ_F can be much higher than TeV since $D(H_u)$ and $D(H_d)$ are close to unity.

A similar operator

$$\frac{F F_u^c F F_d^c}{\Lambda_F} \quad (43)$$

would cause flavor changing neutral current processes. However, if the wave function profiles explain the Yukawa hierarchy, the flavor violation also has the structure of the square of the Yukawa coupling constants. This is a realization of the minimal flavor violation [42, 43, 44].

7 Phenomenology

The properties of the Higgs boson are the same as the Standard Model ones as long as we look at the low energy effective theory. However, there can be contributions from physics at the scale Λ' . Since the Higgs fields couple to charged particles with masses of order $\Lambda' \sim \text{TeV}$, there should be imprints of the strongly coupled sector in the property of the Higgs boson.

The charged particles with masses of order Λ' may give an important contribution to the $h \rightarrow \gamma\gamma$ decay. From the NDA analysis, the effective operator obtained by integrating out these particles are

$$\mathcal{L}_{\text{eff}} \sim \frac{e^2}{(4\pi)^2} \frac{\lambda_d^2 \langle H \rangle}{\Lambda'^2} \cdot h F^{\mu\nu} F_{\mu\nu}. \quad (44)$$

The decay width is then given by

$$\Gamma(h \rightarrow \gamma\gamma) \propto \left| A_{\text{SM}} + \frac{\lambda_d^2 \langle H \rangle^2}{\Lambda'^2} \right|^2, \quad (45)$$

where $A_{\text{SM}} \sim -6.5$ is the contribution from the Standard Model. The deviation from the Standard Model prediction, $\Delta\Gamma/\Gamma_{\text{SM}}$, is therefore estimated to be

$$|\Delta\Gamma/\Gamma_{\text{SM}}| \sim 3\% \cdot \left(\frac{\lambda_d/4\pi}{0.2} \right)^2 \left(\frac{\Lambda'}{1 \text{ TeV}} \right)^{-2}. \quad (46)$$

The sign is not determined. Depending on $O(1)$ parameters, one may be able to see this contribution.

The S , T parameters [45] provide constraints on contributions from the Λ' scale physics. The NDA estimate of the S and T parameters are:

$$\Delta S \sim \frac{1}{\pi} \left(\frac{\lambda_d \langle H \rangle}{\Lambda'} \right)^2 \sim 0.06 \cdot \left(\frac{\lambda_d/4\pi}{0.2} \right)^2 \left(\frac{\Lambda'}{1 \text{ TeV}} \right)^{-2}, \quad (47)$$

$$\begin{aligned} \Delta T &\sim \frac{1}{16\pi s_W^2 c_W^2 m_Z^2} \frac{(\lambda_d \langle H \rangle)^4}{\Lambda'^2} \left(\frac{\langle H_- \rangle}{\langle H_+ \rangle} \right)^2 \\ &\sim 0.09 \cdot \left(\frac{\lambda_d/4\pi}{0.2} \right)^4 \left(\frac{\Lambda'}{1 \text{ TeV}} \right)^{-2} \left(\frac{\langle H_- \rangle / \langle H_+ \rangle}{0.4} \right)^2. \end{aligned} \quad (48)$$

The experimental bounds are $\Delta S, \Delta T \lesssim 0.1$. The S parameter constraint is not very strong. The T parameter is proportional to $\langle H_- \rangle^2$ which represents the violation of the custodial symmetry. By using the estimates in Eqs. (28), (30), and (31), one can see that the constraint from the T parameter is satisfied without any fine-tuning.

At the scale Λ' , there are ρ meson-like vector resonances which have the same quantum numbers as the W and Z bosons. The mixing with the W and Z bosons are expected to be of order, $g_L/(4\pi)$, where g_L is the gauge coupling constant of the $SU(2)_L$ gauge boson. Through this mixing, the resonance can be produced at the LHC via $q\bar{q} \rightarrow W'$ or Z' process. There are various decay modes. For example for the W' boson, it can decay into $t\bar{b}$, $\chi^0\chi^+$, and a pair of a neutral boson (h, Z, H^0) and a charged boson (W, H^+). The branching ratios are weighted by numbers of degrees of freedom and λ^4 . For example,

$$\frac{B(W' \rightarrow WZ)}{B(W' \rightarrow t\bar{b})} \sim \frac{\lambda_d^4}{2\lambda_q^4 N_c}. \quad (49)$$

The most promising decay mode to look for the resonance is $W' \rightarrow WZ$ followed by decays into three leptons. The experimental bound on such a resonance has been studied in Ref. [46]. Due to a small mixing and a small branching fraction, there seems to be no experimental constraint on this resonance for $\Lambda' \sim \text{TeV}$ yet, but there will be a chance to see it in near future.

8 Discussion

The analysis of the two-site model supports that the dynamical electroweak symmetry breaking by QCD is possible in an extra dimensional theory. Not only that, we find that the model is phenomenologically attractive. One can explain the 126 GeV Higgs boson mass without a fine-tuning. These successes give a good motivation to construct a model in string or M-theory. Such a construction will make it completely clear if the emergent Higgs scenario can be embedded to a UV complete higher dimensional theory.

If we extend the model to an N -site model, there should appear a tower of Higgs sectors and techni-rho mesons by sequentially taking the Seiberg dualities. The model then looks like the 5D MSSM where $SU(2)$ gauge fields and the Higgs sector are propagating into an extra dimension. This suggests an interesting duality; a 5D Higgsless MSSM with bulk QCD is dual to another 5D MSSM with the bulk electroweak sector including the Higgs fields.

The two-site model we studied has the same structure as the one in Ref. [47], where it is tempted to understand the ρ meson in QCD as the magnetic gauge boson. This

exhibits a similarity of the chiral symmetry breaking in QCD and the electroweak symmetry breaking. Our construction of the extra dimensional model provides a unification of these two phenomena.

SUSY is probably essential for defining a higher dimensional gauge theory. Not only for a theory definition, SUSY explains why the Higgs boson is so light. However, in our discussion, a light Higgs boson is obtained by a small coupling constant rather than a low SUSY breaking scale. From the bottom up point of view, the point is the presence of a UV completion in which the Higgs operators do not have a potential at tree level, and couple to some TeV dynamics with the strength of $\lambda/4\pi \sim 0.2$ rather than unity. We are not aware of, but there can be such examples of non-supersymmetric models.

Whether or not the Higgs fields are fundamental is probably not a physical question. Once we have a theory, they are just effective degrees of freedom wherever they come from. An important message from this study is that a trial to define a higher dimensional theory leads us to a picture of emergent Higgs fields. A generic prediction of the framework is then a presence of relatively light resonances at a scale $\Lambda' \sim 1$ TeV. If the Higgs boson is weakly coupled as suggested by its mass, the techni-rho mesons should be lighter than the naive estimate, $\Lambda' < 4\pi\langle H \rangle \sim 2 - 3$ TeV. We expect to see them at the LHC.

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